Z-mediated charge and CP asymmetries and FCNCs in $B_{d,s}$ processes

Chuan-Hung Chen[1,2a], **Chao-Qiang Geng**[3b] **and Wei Wang**[4c]

[1] *Department of Physics, National Cheng-Kung University, Tainan 701, Taiwan*

[2] *National Center for Theoretical Sciences, Hsinchu 300, Taiwan*

[3] *Department of Physics, National Tsing-Hua University, Hsinchu 300, Taiwan*

[4] *Istituto Nazionale di Fisica Nucleare, Sezione di Bari, Bari 70126, Italy*

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Abstract

We show model-independently that the negative like-sign charge asymmetry ($-A^b_{s\ell}$) is less than 3.16×10^{-3} when the constraints from the $B_q-\bar{B}_q$ mixings and the time-dependent CP asymmetries (CPAs) for $B_q \to J/\Psi M_q$ with $M_q = K, \phi$ and $q = d, s$ are taken into account. Although the result is smaller than the measured value by the DØ Collaboration at Fermilab, there is still plenty of room to have new physics, which is sensitive to new CP violating effects, as the standard model (SM) prediction is $(2.3^{+0.6}_{-0.5}) \times 10^{-4}$. To illustrate the potential large $|A^b_{s\ell}|$, we show the influence of new $SU(2)_L$ singlet exotic quarks in the vector-like quark model, where the Z-mediated flavor changing neutral currents (FCNCs) are generated at tree level. In particular, we demonstrate that (a) the like-sign charge asymmetry could be enhanced by a factor of two in magnitude; (b) the CPA of $\sin 2\beta_s^{J/\Psi\phi}$ could reach to -15%; (c) the CPA of $\sin 2\beta_{\phi K_S}$ could be higher than $\sin 2\beta_{J/\Psi K_S}$ when $|A^b_{s\ell}|$ is larger than the SM prediction; and (d) the branching ratio for $B_s \to \mu^+\mu^-$ could be as large as 0.6×10^{-8}.

[a] Email: physchen@mail.ncku.edu.tw

[b] Email: geng@phys.nthu.edu.tw

[c] Email: wei.wang@ba.infn.it

I. INTRODUCTION

It is clear that some new CP violation mechanism beyond the Kobayashi-Maskawa (KM) phase in the standard model (SM) is needed in order to explain the matter-antimatter asymmetry of the Universe. Moreover, several hints for the existence of some new CP violating phases are revealed in the low energy processes, such as the πK puzzle in $B \to \pi K$ decays, the large CP asymmetry (CPA) of $\sin 2\beta_s^{J/\Psi\phi}$ in the $B_s \to J/\Psi\phi$ decay, inconsistent time-dependent CPAs between $B_d \to (\eta, \phi) K_S$ and $B_d \to J/\Psi K_S$ decays, etc [1].

Recently, the DØ Collaboration at Fermilab has observed the like-sign charge asymmetry, defined as [2]

$$A^b_{s\ell} = \frac{N_b^{++} - N_b^{--}}{N_b^{++} + N_b^{--}}\,, \tag{1}$$

where $N_b^{++(--)}$ denotes the number of events that b and $\bar{b}$-hadron semileptonically decay into two positive (negative) muons. The measured value in the dimuon events is given by [2]

$$A^b_{s\ell} = (-9.57 \pm 2.51(\mathrm{stat}) \pm 1.46(\mathrm{syst})) \times 10^{-3}\,, \tag{2}$$

which is about 3.2 standard deviations from the SM prediction of $(-2.3^{+0.5}_{-0.6}) \times 10^{-4}$ [2, 3]. If the semileptonic b-hadron decays do not involve a CP phase, the charge asymmetry is directly related to the mixing-induced CPAs in $B_{d,s}$-meson oscillations (see the detailed analysis later). Although the errors of the data are still large, the deviations from the SM could be attributed to the new CP violating phases in $b \to d$ and $b \to s$ transitions [4–16].

Inspired by the new DØ measurement and other CPAs measured earlier, we illustrate that the anomalies can be induced by the new exotic vector-like quarks in the so-called vector-like-quark model (VQM). Unlike the conventional four-generation model with the fourth left-handed quarks being an $SU(2)_L$ doublet, the vector-like quarks (VQs) are all $SU(2)_L$ singlets, as the ones naturally realized in E_6 models [17]. Since the left-handed VQs carry the same hypercharge as the right-handed quarks in the SM, interestingly the model leads to Z-mediated flavor changing neutral currents (FCNCs) at tree level [18–21]. Moreover, the VQM involves less free parameters and is more predictable since the couplings of Z-boson to fermions and m_Z are known. In addition to the mixing-induced CPAs, the VQM has significant impacts on the rare B_q decays such as $b \to s\ell^+\ell^-$ and $B_s \to \mu^+\mu^-$ as well as other B_q processes.

The paper is organized as follows. In Sec. II, we analyze model-independently the wrong and like-sign charge asymmetries in detail. In Sec. III, we derive Feynman rules for the Z-mediated FCNCs in the VQM and formulate CPAs and rare B_q decays. The numerical analysis is presented in Sec. IV. The conclusion is given in Sec. V

II. MODEL-INDEPENDENT RESULTS ON CHARGE ASYMMETRIES

In order to comprehend the implication of the like-sign charge asymmetry $A^b_{s\ell}$, we use both experimental and phenomenological approaches. We first discuss the issue from the viewpoint of the current data. To evaluate $A^b_{s\ell}$, we start with the wrong-sign charge asymmetry in semileptonic B_q decays, defined by [25]

$$\begin{aligned} a^q_{s\ell} &= \frac{\Gamma(\bar{B}_q(t) \to \ell^+ X) - \Gamma(B_q(t) \to \ell^- X)}{\Gamma(\bar{B}_q(t) \to \ell^+ X) + \Gamma(B_q(t) \to \ell^- X)}\,, \\ &\approx Im\left(\frac{\Gamma^q_{12}}{M^q_{12}}\right) \end{aligned} \tag{3}$$

where $\Gamma^q_{12}(M^q_{12})$ denotes the absorptive (dissipative) part of the $B_q \leftrightarrow \bar{B}_q$ transition with $\Gamma^q_{12} \ll M^q_{12}$. As a consequence, a non-zero $a^q_{s\ell}$ indicates CP violation. As Γ^q_{12} is dominated by the SM contributions, we adopt $\Gamma^q_{12} = \Gamma^q_{12}(\text{SM})$ in the following analysis. The SM predictions are $a^d_{s\ell}(\text{SM}) = (-4.8^{+1.0}_{-1.2}) \times 10^{-4}$ and $a^s_{s\ell}(\text{SM}) = (2.06 \pm 0.57) \times 10^{-5}$ [3], while the current data are $a^d_{s\ell}(\text{Exp}) = (-4.7 \pm 4.6) \times 10^{-3}$ [1] and $a^s_{s\ell}(\text{Exp}) = (-1.7 \pm 9.1) \times 10^{-3}$ [24]. The relation between the wrong and like-sign charge asymmetries indeed can be expressed by [2, 23]

$$\begin{aligned} A^b_{s\ell} &= \frac{\Gamma(b\bar{b} \to \ell^+\ell^+ X) - \Gamma(b\bar{b} \to \ell^-\ell^- X)}{\Gamma(b\bar{b} \to \ell^+\ell^+ X) + \Gamma(b\bar{b} \to \ell^-\ell^- X)}\,, \\ &= 0.506(43) a^d_{s\ell} + 0.494(43) a^s_{s\ell}\,. \end{aligned} \tag{4}$$

From Eq. (4), it is easy to see that the like-sign charge asymmetry depends on the CP phases in B_d and B_s oscillations. If we take $a^d_{s\ell}(\text{Exp})$ and the DØ observed value of $A^b_{s\ell}$ as inputs, we immediately get

$$A^s_{s\ell} = 0.494(43) a^s_{s\ell} = (-7.2 \pm 3.7) \times 10^{-3}\,. \tag{5}$$

In other words, the wrong-sign charge asymmetry $a^s_{s\ell}$ can be extracted as

$$a^s_{s\ell}(\text{Extr}) = -0.01456 \pm 0.00764\,, \tag{6}$$

where the errors have been regarded as uncorrelated and combined in quadrature. Similarly, if $a^d_{s\ell}$ is negligible, $a^s_{s\ell}(\text{Extr}) = -0.01937 \pm 0.0061$. Clearly, by the current experimental values, $|a^s_{s\ell}(\text{Extr})|$ is three orders of magnitude larger than the SM prediction.

After discussing the allowed value of $a^q_{s\ell}$ from the viewpoint of the current experimental data, it is interesting to analyze the same wrong-sign charge asymmetry from Eq. (3) directly. We set $\Gamma^q_{12} = \Gamma^q_{12}(\text{SM}) = -|\Gamma^q_{12}(\text{SM})|e^{i\phi^\Gamma_q}$ and write the $B_q - \bar{B}_q$ transition matrix element as

$$\begin{aligned} M^q_{12} &= M^q_{12}(SM) + M^q_{12}(NP)\,, \\ &= |M^q_{12}(SM)| R_q \exp(2i\beta_q + i\phi^{NP}_q) \end{aligned} \tag{7}$$

where

$$\begin{aligned} R_q &= \left(1 + r_q^2 + 2r_q \cos 2(\theta^{NP}_q - \beta_q)\right)^{1/2}\,, \\ r_q &= \frac{|M^q_{12}(NP)|}{|M^q_{12}(SM)|}\,, \\ 2\beta_q &= \arg(M^q_{12}(SM))\,,\ 2\theta^{NP}_q = \arg(M^q_{12}(NP))\,, \\ \tan\phi^{NP}_q &= \frac{r_q \sin 2(\theta^{NP}_q - \beta_q)}{1 + r_q \cos 2(\theta^{NP}_q - \beta_q)}\,. \end{aligned} \tag{8}$$

With $\Delta\Gamma^q = 2|\Gamma^q_{12}|\cos\phi_q$ and $\phi_q = arg\,(-M^q_{12}/\Gamma^q_{12})$, Eq. (3) could be expressed as

$$a^q_{s\ell} \approx Im\left(\frac{\Gamma^q_{12}}{M^q_{12}}\right) \approx \frac{\Delta\Gamma^q(\text{SM})}{\Delta m_{B_q}} \frac{\sin\left(\phi^{\text{NP}}_q + \phi^{\text{SM}}_q\right)}{\cos\phi^{\text{SM}}_q}\,. \tag{9}$$

By using the SM results [3]:

$$\begin{aligned} \Delta\Gamma_d(\text{SM}) &= (2.67^{+0.58}_{-0.65}) \times 10^{-3}\ \text{ps}^{-1}\,, \\ \Delta\Gamma_s(\text{SM}) &= 0.096 \pm 0.039\ \text{ps}^{-1}\,, \\ \phi^{\text{SM}}_d &= -0.091^{+0.026}_{-0.038}\,, \\ \phi^{\text{SM}}_s &= (4.3 \pm 1.4) \times 10^{-3}\,, \end{aligned} \tag{10}$$

and the data: $\Delta m_{B_d} = 0.508 \pm 0.005\ \text{ps}^{-1}$ and $\Delta m_{B_s} = 17.77 \pm 0.12\ \text{ps}^{-1}$ [25], we obtain

$$\begin{aligned} a^d_{s\ell} &= (5.26^{+1.14}_{-1.28}) \times 10^{-3} \sin(\phi^{\text{NP}}_d + \phi^{\text{SM}}_d)\,, \\ a^s_{s\ell} &= (5.40 \pm 2.20) \times 10^{-3} \sin(\phi^{\text{NP}}_s + \phi^{\text{SM}}_s)\,. \end{aligned} \tag{11}$$

Obviously, the sign and magnitude of $a^q_{s\ell}$ are dictated by the factor of $\sin\left(\phi^{\text{NP}}_q + \phi^{\text{SM}}_q\right)$. We find that the most strict model-independent constraints on $\sin(\phi^{\text{NP}}_q + \phi^{\text{SM}}_q)$ are from Δm_{B_q} and the time-dependent CPA of $S_{J/\Psi M_q} = \sin(2\beta_q + \phi^{\text{NP}}_q)$ [7, 25] for $B_q \to J/\Psi M_q$ with

$M_q = K(\phi)$ and $q = d(s)$. We note that since $B_q \to J/\Psi M_q$ is dominated by tree diagrams in the SM, we have assumed that the contribution to the decay amplitude from new physics is negligible. With $\Delta m_{B_d}(\text{SM}) = 0.506\ \text{ps}^{-1}$, $\Delta m_{B_s}(\text{SM}) = 17.80\ \text{ps}^{-1}$, $\beta_d = 0.38 \pm 0.01$ [26], $\beta_s \approx -0.019$ [19], $S^{\text{Exp}}_{J/\Psi K_S} = 0.655 \pm 0.024$, $S^{\text{Exp}}_{J/\psi\phi} \in (-0.995, -0.285)$ [1] and $(\Delta m_{B_q})^{\text{Exp}}$, we derive

$$
\begin{aligned}
-\sin(\phi^{\text{NP}}_d + \phi^{\text{SM}}_d) &< 0.2\,, \\
-\sin(\phi^{\text{NP}}_s + \phi^{\text{SM}}_s) &< 0.985\,.
\end{aligned} \tag{12}
$$

In Figs. 1(a) and 2(a), we present Δm_{B_q} and $S_{J/\Psi M}$ with 2σ errors of the data as functions of r_q and θ^{NP}_q, while the contours for $\sin(\phi^{\text{NP}}_q + \phi^{\text{SM}}_q)$ are displayed in Figs. 1(b) and 2(b) for $q = d$ and s, respectively. In Fig. 1(b), the scattered patten denotes the combined constraints from data of Δm_{B_d} and $S_{J/\Psi K_S}$. If we take the central values in Eq. (11) as inputs, we immediately obtain that $-a^d_{s\ell} < 1.05 \times 10^{-3}$ and $-a^s_{s\ell} < 5.32 \times 10^{-3}$. Although the central values of $|a^d_{s\ell}(\text{Exp})|$ and $|a^s_{s\ell}(\text{Extr})|$ are larger than our phenomenological analysis, both results are still consistent with each other when the errors of the data are taken into account. As a result, we obtain the model-independent (MI) result on the negative like-sign

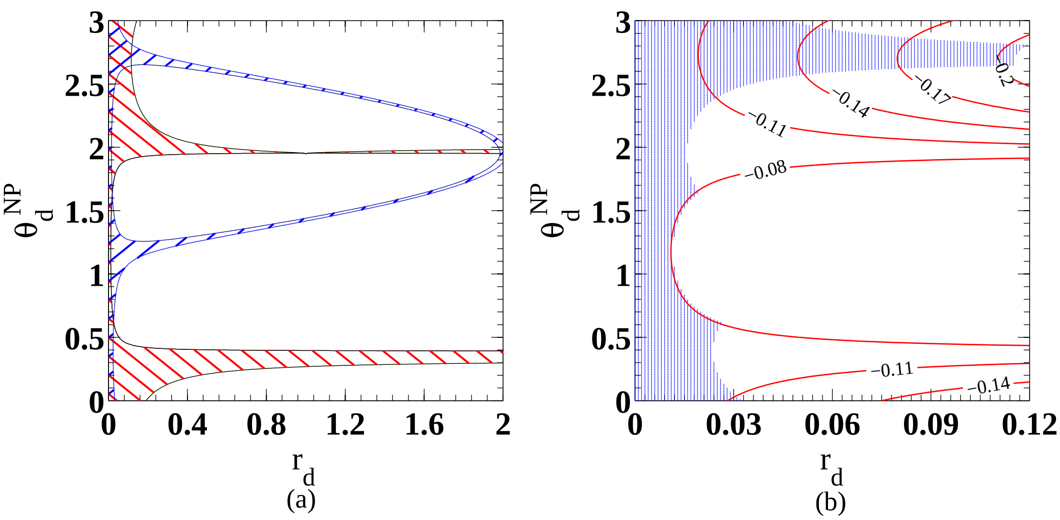

FIG. 1. (a) Constraints from 2σ errors of $(\Delta m_{B_d})^{\text{Exp}}$ (down-left hatched) and $S^{\text{Exp}}_{J/\Psi K_S}$ (down-right hatched) and (b) Contours for $\sin(\phi^{\text{NP}}_d + \phi^{\text{SM}}_d)$ as a function of r_d and θ^{NP}_d.

charge asymmetry, given by

$$
-A^b_{s\ell}(MI) = -0.506(43)a^d_{s\ell}(MI) - 0.494(43)a^s_{s\ell}(MI) < 3.16 \times 10^{-3}\,. \tag{13}
$$

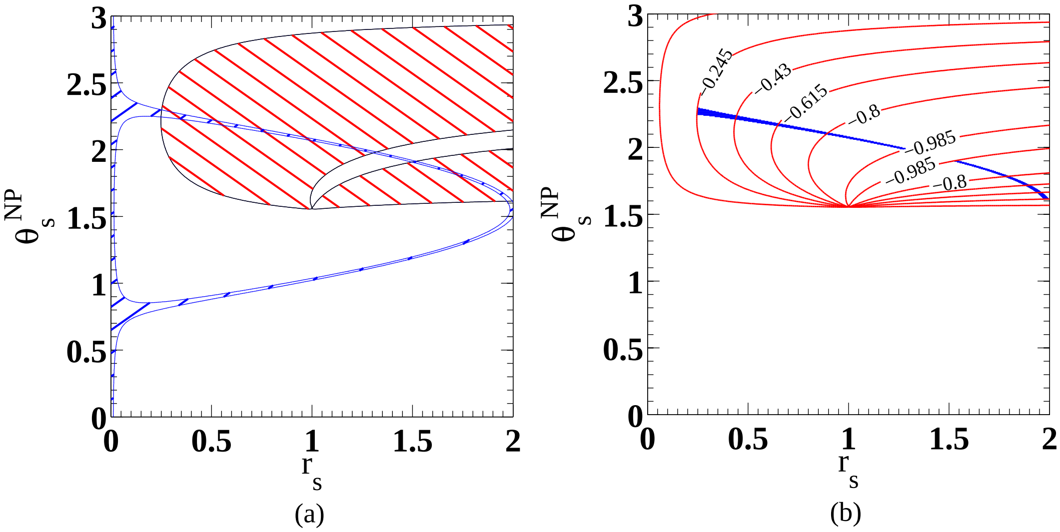

FIG. 2. Legend is the same as Fig. 1 but for q=s.

Although the value in Eq. (13) is smaller than the measured value by DØ in Eq. (2), it is still one order of magnitude larger than the SM prediction. Therefore, $A^b_{s\ell}$ could be a good candidate to probe the new CP violating source in the $B_{d,s}$ systems at Tevatron, LHCb and super-B factories.

III. VECTOR-LIKE QUARK MODEL

A. Z-mediated FCNCs

By extending the SM with including the new $SU(2)_L$ singlet down quarks of D_L and D_R, the extended Yukawa sector becomes

$$-\mathcal{L}_Y = \bar{Q}_L Y_D H d_R + h_D \bar{Q}_L H D_R + m_D \bar{D}_L D_R + h.c.\,, \tag{14}$$

where we have suppressed the flavor indices, Q_L (H) is the SU(2) quark (Higgs) doublet, Y_D and h_D are Yukawa couplings and m_D is the mass of the exotic quark before the electroweak symmetry breaking. When the Higgs field develops the vacuum expectation value (VEV), the mass matrix of the down type quark is given by

$$m_d = \left(\begin{array}{c|c} Y^{ij}_D & 0 \\ \hline h^j_D & m_D \end{array} \right) . \tag{15}$$

Introducing two unitary matrices, the mass matrix can be diagonalized by

$$m_d^{\rm dia} = V_D^L m_d V_D^{R\dagger} \,. \tag{16}$$

In the SM, since the interactions of Z-boson to fermions are flavor blind, the flavor in the process with the exchange of Z-boson is naturally conserved at tree level. In the VQM, the new left-handed quark is an $SU(2)_L$ singlet and carries the same hypercharge as the right-handed down-type quarks. The gauge interactions of the left-handed down-type quarks with Z-boson are given by

$$\begin{aligned} \mathcal{L}_Z &= -\frac{g c_L^f}{2\cos\theta_W} \bar{F}\gamma^\mu X_F P_L F Z_\mu \,, \\ X_F &= \left[\begin{array}{c|c} \mathbb{1}_{3\times 3} & \mathbf{0}_{3\times 1} \\ \hline \mathbf{0}_{1\times 3} & \xi_D \end{array} \right] , \end{aligned} \tag{17}$$

where g is the coupling constant of $SU(2)_L$, θ_W is the Weinberg's angle, $P_{R(L)} = (1 \pm \gamma_5)/2$, $F^T = (d, s, b, D)$ represents the down-type quarks including the new singlet, c_L^f is defined as $c_{L(R)}^f = c_V^f \pm c_A^f$ with

$$c_V^f = I_f^3 - 2\sin^2\theta_W Q_f \,, \quad c_A^f = I_f^3 \tag{18}$$

in which I_f^3 and Q_f are the third component of the weak isospin and the electric charge of the particle, respectively, and $\xi_f = -2\sin^2\theta_W Q_f / c_L^f$. Due to $X_F \neq \mathbb{1}_{4\times 4}$, accordingly, Eq. (17) leads to FCNCs at tree level. Since D_R and $q_R = (d, s, b)_R$ have the same quantum number, the right-handed quarks are FCNC free at tree level. Following Eq. (16), the couplings of Z-boson to fermions in the mass eigenstates are written by

$$\mathcal{L}_Z = -\frac{g c_L^f}{2\cos\theta_W} \bar{F}\gamma^\mu \left(V_D^L X_F V_D^{L^\dagger} \right) P_L F Z_\mu \,. \tag{19}$$

The FCNC effects could be further formulated as

$$\left(V_D^L X_F V_D^{L^\dagger} \right)_{f'f} = \delta_{f'f} + (V_D^L)_{f'D} (\xi_D - 1) (V_D^{L^*})_{fD} = \delta_{f'f} + \lambda_{f'f} \,. \tag{20}$$

Thus, the interaction for b-q-Z is given by

$$\mathcal{L}_{b\to q} = -\frac{g c_L^d \lambda_{qb}}{2\cos\theta_W} \bar{q}\gamma^\mu P_L b Z_\mu + h.c. \tag{21}$$

with

$$\lambda_{qb} = (\xi_D - 1)(V_D^L)_{qD} (V_D^L)^*_{bD} \equiv |\lambda_{qb}| \exp\left(i\theta_q^Z \right) \,.$$

Clearly, the new free parameters are only λ_{db} and λ_{sb}. When λ_{qb} is fixed by the current data, one may have some solid predictions for the relevant processes.

B. $B_q - \bar{B}_q$ mixing

With Eq. (21) and the hadronic transition matrix element defined by

$$\langle B_q|\bar{q}\gamma_\mu P_{L(R)} b\bar{q}\gamma_\mu P_{L(R)} b|\bar{B}_q\rangle = \frac{1}{3} m_{B_q} f^2_{B_q} \hat{B}_q \,, \tag{22}$$

the matrix element for $\bar{B}_q \to B_q$ mediated by the Z-boson at tree level is obtained as

$$M^q_{12}(Z) = \frac{G_F \left(\lambda_{qb} c^d_L\right)^2}{3\sqrt{2}} m_{B_q} f^2_{B_q} \hat{B}_q = |M^q_{12}(Z)| e^{2i\theta^Z_q} \,. \tag{23}$$

In addition to the tree effects, the Z-mediated box and penguin diagrams will induce important linear term in λ_{qb} and it is given by [21]

$$M^q_{12}(Loop) = -1.3 \lambda_{qb} V^*_{tq} V_{tb} \,. \tag{24}$$

Following Eq. (7), the combination of the SM and Z-mediated tree, box and penguin contributions for the $B_q - \bar{B}_q$ mixing is given by

$$\begin{aligned} M^q_{12} &= M^q_{12}(SM) + M^q_{12}(Z) + M^q_{12}(Loop) \\ &= M^q_{12}(SM) R^Z_q e^{2i\beta^Z_q} \,, \end{aligned} \tag{25}$$

where the corresponding parameters in Eq. (8) could be obtained by the following replacements: $M^q_{12}(Z) + M^q_{12}(Loop) = M^q_{12}(NP)$, $R^Z_q = R_q(r^Z_q, \theta^Z_q)$ and $\phi^{NP}_q = \phi^Z_q = 2\beta^Z_q$. Hence, the mixing parameter for the B_q oscillation is $\Delta m_{B_q} = 2|M^q_{12}(SM)| R^Z_q = R^Z_q \Delta m_{B_q}(SM)$.

C. Mixing-induced CP asymmetries

After deriving M^q_{12}, we now can study the mixing-induced CPAs. The first types of CPAs are the wrong and like-sign charged asymmetries, defined in Eqs. (3) and (4). Since the relationship between the wrong and like-sign asymmetries has been given in Eq. (4), we simply formulate the Z-mediated $a^q_{s\ell}$ as

$$\begin{aligned} a^q_{s\ell} &= Im\left(\frac{\Gamma^q_{12}}{M^q_{12}}\right) \,, \\ &\approx \frac{\Delta\Gamma^q(SM)}{\Delta m_{B_q}(SM)\cos\phi^{SM}_q} \frac{\sin(\phi^{SM}_q + \phi^Z_q)}{R^Z_q} \,, \end{aligned} \tag{26}$$

where all SM related quantities are taken to be known. Note that $a^q_{s\ell}$ involves two free parameters, i.e. $|\lambda_{qb}|$ and θ^Z_q.

Another type of the time-dependent CPA is associated with the definite CP in the final state, defined by [25]

$$
\begin{aligned}
A_{f_{CP}}(t) &= \frac{\Gamma(\bar{B}_q(t) \to f_{CP}) - \Gamma(B_q(t) \to f_{CP})}{\Gamma(\bar{B}_q(t) \to f_{CP}) + \Gamma(B_q(t) \to f_{CP})}\,, \\
&= S_{f_{CP}} \sin \Delta m_{B_q} t - C_{f_{CP}} \cos \Delta m_{B_q} t\,, \\
S_{f_{CP}} &= \frac{2Im\lambda_{f_{CP}}}{1+|\lambda_{f_{CP}}|^2}\,, \quad C_{f_{CP}} = \frac{1-|\lambda_{f_{CP}}|^2}{1+|\lambda_{f_{CP}}|^2}\,,
\end{aligned} \tag{27}
$$

with

$$
\lambda_{f_{CP}} = -\left(\frac{M_{12}^{B_q^*}}{M_{12}^{B_q}}\right)^{1/2} \frac{A(\bar{B} \to f_{CP})}{A(B \to f_{CP})} = -e^{-2i(\beta_q + \phi_q^{\rm NP})} \frac{\bar{A}_{f_{CP}}}{A_{f_{CP}}}\,, \tag{28}
$$

where f_{CP} denotes the final CP eigenstate, $S_{f_{CP}}$ and $C_{f_{CP}}$ are the so-called mixing-induced and direct CPAs, $A_{f_{CP}}$ and $\bar{A}_{f_{CP}}$ are the amplitudes of B and $\bar{B}$ mesons decaying to f_{CP} and $\bar{A}_{f_{CP}}/A_{f_{CP}} = -\eta_{f_{CP}} A_{f_{CP}}(\theta_W \to -\theta_W)/A_{f_{CP}}(\theta_W)$ with $\eta_{f_{CP}}$ and θ_W being the CP eigenvalue of f_{CP} and the weak CP phase, respectively. Clearly, besides the phase in the $\Delta B = 2$ process, the mixing-induced CPA is also related to the phase in the $\Delta B = 1$ process. Due to $B \to \eta' K_S$ involving more complicated and uncertain QCD effects, in this paper, we will concentrate on $f_{CP} = J/\Psi K_S$ and ϕK_S for $q = d$ and $f_{CP} = J/\Psi \phi$ for $q = s$.

For $\Delta B = 1$ processes, we also need to know the flavor conserving interactions. The couplings of Z-boson to fermions in the SM are summarized as

$$
\mathcal{L}_Z^{SM} = -\frac{g}{2\cos\theta_W} \sum_f \bar{f}\gamma^\mu \left(c_V^f - c_A^f \gamma_5\right) f Z_\mu\,, \tag{29}
$$

where f denotes any fermions and $c^f_{V(A)}$ is given in Eq. (18). Using Eqs. (21) and (29), the Z-mediated Hamiltonian for $b \to q q' \bar{q}'$ decays is obtained by

$$
\mathcal{H}^Z_{b\to qq'\bar{q}'} = \frac{G_F}{\sqrt{2}} \left(\frac{\lambda_{qb} c_L^d}{2}\right) (\bar{q}b)_{V-A} \sum_{q'=u,d,s,c} \left(c_L^{q'} (\bar{q}'q')_{V-A} + c_R^{q'} (\bar{q}'q')_{V+A}\right) \tag{30}
$$

where $(\bar{f}'f)_{V\pm A} = \bar{f}'\gamma^\mu(1\pm\gamma_5)f$. Clearly, the Z-mediated effects for $b \to qq'\bar{q}'$ are similar to the standard electroweak penguins but the SM contributions are small. Since $B_d \to J/\Psi K_S$ and $B_s \to J/\Psi \phi$ decays are dominated by the tree diagrams, the penguin-like effects can be regarded to be relatively small and insignificant in the $b \to sc\bar{c}$ processes. On the contrary, since $b \to ss\bar{s}$ is a penguin dominant process, the Z-mediated effects are naturally comparable with the SM contributions. Hence, we will only focus on $B \to \phi K_S$. In order to deal with the hadronic effects in nonleptonic B_q decays, we employ the naive factorization

approach (NFA). The decay amplitude combined the SM with Z-mediated contributions for $B \to \phi K$ is written as

$$\begin{aligned}
\bar{A}_{\phi \bar{K}^0} &= \langle \phi \bar{K} | \mathcal{H}_{b\to ss\bar{s}} | \bar{B}^0 \rangle \,, \\
&= \frac{G_F}{\sqrt{2}} V_{ts}^* V_{tb} (a^{\rm SM} + a_s^Z) \langle \phi | \bar{s}\gamma_\mu s | 0 \rangle \langle \bar{K}^0 | \bar{s}\gamma^\mu b | \bar{B} \rangle \,, \qquad (31)
\end{aligned}$$

with

$$\begin{aligned}
a_s^{\rm SM} &= a_3 + a_4 + a_5 \,, \\
a_3 &= C_3 + \frac{C_4}{N_C} \,, \quad a_4 = C_4 + \frac{C_3}{N_C} \,, \\
a_5 &= C_5 + \frac{C_6}{N_C} \,, \quad a_s^Z = -\frac{\lambda_{sb} c_L^d}{V_{ts}^* V_{tb}} \left(c_V^s + \frac{c_L^s}{2N_C} \right) \,,
\end{aligned}$$

where N_C is the number of colors and C_{3-6} are the effective Wilson coefficients from the gluon penguins of the SM [28]. We note that the electroweak penguin contributions are very small and neglected in the analysis.

Consequently, the ratio of amplitudes for $\bar{B}_d \to \phi K_S$ and $B_d \to \phi K_S$ decays is written as

$$\frac{\bar{A}_{\phi K_S}}{A_{\phi K_S}} = -e^{2i\beta_s} \frac{a^{SM} + a_s^Z}{a^{SM} + a_s^{Z*}} = -e^{2i(\beta_s + \delta_s^Z)} \qquad (32)$$

with

$$\tan \delta_s^Z = \frac{|a_s^Z| \sin(\theta_s^Z - \beta_s)}{a^{SM} + |a_s^Z| \cos(\theta_s^Z - \beta_s)} \,.$$

From Eqs. (27) and (28), the mixing-induced CPA through the ϕK_S mode is obtained as

$$S_{\phi K_S} \equiv \sin 2\beta_{\phi K_S} = \sin 2(\beta_d + \beta_d^Z - \beta_s - \delta_s^Z) \,. \qquad (33)$$

Similarly, the CPAs through $J/\Psi(K_S, \phi)$ channels are simply given by

$$\begin{aligned}
S_{J/\Psi K_S} &\equiv \sin 2\beta_{J/\Psi K_S} \approx \sin(2\beta_d + \phi_d^Z) \,, \\
S_{J/\Psi \phi} &\equiv \sin 2\beta_s^{J/\Psi\phi} \approx \sin(2\beta_s + \phi_s^Z) \,. \qquad (34)
\end{aligned}$$

Although $\sin 2\beta_{J/\Psi K_S}$ has been measured at a precision level, it might be difficult to confirm whether new physics exists by observing $\sin 2\beta_{J/\Psi K_S}$ alone. Nevertheless, one can investigate a new asymmetry, defined by [30]

$$\Delta S_{\beta_d} = \sin 2\beta_{J/\Psi K_S} - \sin 2\beta_{\phi K_S} \,, \qquad (35)$$

in which the SM prediction is less than around 5% [30]. Clearly, if a large value of ΔS_{β_d} is measured, it will be a strong hint for new physics beyond the SM.

D. $b \to q\ell^+\ell^-$ and $B_q \to \mu^+\mu^-$ decays

In addition to the CP violating observables, the other interesting environment to probe the new physics effects is rare decays in which the predicted branching ratios (BRs) in the SM are small. Although the BR is not a direct CP violating observable, it is still sensitive to the CP violating effect via the squared imaginary coupling. In most exclusive decay processes, the BRs are associated with uncertain nonperturbative hadronic effects. To reduce the QCD uncertainties, we choose inclusive $b \to q\ell^+\ell^-$ and exclusive $B_q \to \ell^+\ell^-$ decays as the candidates to probe the new physics effects, where the hadronic effects could be controlled well.

Using Eqs. (21) and (29), the effective Hamiltonian for $b \to q\ell^+\ell^-$ mediated by Z-boson is found to be

$$\mathcal{H}^Z_{b\to q\ell^+\ell^-} = \frac{G_F}{\sqrt{2}}\lambda_{qb}c^d_L(\bar{q}b)_{V-A}\left[c^\ell_V(\bar{\ell}\ell)_V - c^\ell_A(\bar{\ell}\ell)_A\right] \tag{36}$$

where $(\bar{\ell}\ell)_V = \bar{\ell}\gamma^\mu\ell$ and $(\bar{\ell}\ell)_A = \bar{\ell}\gamma^\mu\gamma_5\ell$. Combining with the SM contributions, the decay amplitude for $b \to q\ell^+\ell^-$ is written as

$$\begin{aligned}\mathcal{H}_{b\to q\ell^+\ell^-} = &-\frac{G_F\alpha}{\sqrt{2}\pi}V^*_{tq}V_{tb}\left[\left(C_9\bar{q}\gamma_\mu P_L b - \frac{2m_b}{k^2}C^{SM}_{7\gamma}\bar{q}i\sigma_{\mu\nu}k^\nu P_R b\right)\bar{\ell}\gamma^\mu\ell\right.\\ &\left.+\, C_{10}\bar{q}\gamma_\mu P_L b\bar{\ell}\gamma_\mu\gamma_5\ell\right]\end{aligned} \tag{37}$$

where $k_\mu = (p_{\ell^+} + p_{\ell^-})_\mu$, k^2 is the invariant mass of the lepton pair, and

$$\begin{aligned}C_9 &= C^{SM}_9(m_b) - \frac{2\pi}{\alpha}\frac{\lambda_{qb}c^d_L c^\ell_V}{V^*_{tq}V_{tb}}\,,\\ C_{10} &= C^{SM}_{10} + \frac{2\pi}{\alpha}\frac{\lambda_{qb}c^d_L c^\ell_A}{V^*_{tq}V_{tb}}\,.\end{aligned} \tag{38}$$

The explicit expressions of $C^{SM}_{9,10}$ could be found in Ref. [28]. Accordingly, the differential decay rate is [28]

$$\begin{aligned}\frac{d\Gamma(b\to q\ell^+\ell^-)}{d\hat{s}} &= \Gamma(b\to ce\bar{\nu}_e)\frac{|V^*_{ts}|^2}{|V_{cb}|^2}\frac{\alpha^2}{4\pi^2}\frac{(1-\hat{s})^2}{f(z)k(z)}\\ &\times\left[(1+2\hat{s})\left(|C_9|^2+|C_{10}|^2\right)+4\left(1+\frac{2}{\hat{s}}\right)|C^{SM}_{7\gamma}|^2 + 12C^{SM}_7 ReC_9\right]\,,\\ f(z) &= 1-8z^2+8z^6-z^8-24z^4\ln z\,,\\ k(z) &= 1-\frac{2\alpha_s}{3\pi}\left[\left(\pi^2-\frac{31}{4}\right)(1-z)^2+\frac{3}{2}\right]\,,\end{aligned} \tag{39}$$

where $\hat{s} = k^2/m_b^2$, $z = m_c/m_b$ and $\Gamma(b \to c e \bar{\nu}_e)$ is used to cancel the uncertainties from the CKM matrix elements and m_b^5. Moreover, with Eq. (37) and the B_q meson decay constant, defined by

$$\langle 0|\bar{q}\gamma^\mu\gamma_5 b|\bar{B}_q\rangle = i f_{B_q} p^\mu_{B_q} \,, \tag{40}$$

the BR for $B_q \to \ell^+\ell^-$ is straightforwardly obtained by

$$\mathcal{B}(B_q \to \ell^+\ell^-) = \mathcal{B}^{SM}(B_q \to \ell^+\ell^-) \left| 1 - \frac{\pi}{\alpha} \frac{\lambda_{qb} c^d_L}{V^*_{tq} V_{tb} C^{SM}_{10}} \right|^2 , \tag{41}$$

where

$$\mathcal{B}^{SM}(B_q \to \ell^+\ell^-) = \tau_{B_q} \frac{G_F^2 \alpha^2}{16\pi^3} |V^*_{tq} V_{tb}|^2 m_{B_q} f^2_{B_q} m^2_\ell |C^{SM}_{10}|^2 \left(1 - \frac{4m^2_\ell}{m^2_{B_q}}\right)^{1/2} .$$

IV. NUMERICAL ANALYSIS

As stated earlier, there are four new unknown parameters for $b \to (d, s)$ transitions in the VQM, i.e., $|\lambda_{db,sb}|$ and $\theta^Z_{d,s}$. Although we have model-independently shown the possible severe constraints in Sec. II, in a specific model, we have to consider more relevant bounds. As $\lambda_{f'f} = (\xi_D - 1)(V^L_D)^*_{f'D}(V^L_D)_{fD}$ defined in Eq. (20), the $s \to d$ transition is associated with $(V^L_D)^*_{dD}(V^L_D)_{sD}$ while the $b \to (d, s)$ ones depend on $(V^L_D)^*_{dD}(V^L_D)_{bD}$ and $(V^L_D)^*_{sD}(V^L_D)_{bD}$, respectively. Thus, $(V^L_D)^*_{dD}$ and $(V^L_D)^*_{sD}$ appearing in $b \to (d, s)$ also occur in $s \to d$. We see clearly that $K^0 - \bar{K}^0$ and $B_{d,s} - \bar{B}_{d,s}$ mixings are strongly correlated. Since Δm_K and the indirect (direct) CP violating parameters denoted by ϵ_K (ϵ'_K) are much smaller than those in the B_q systems, the stringent constraints could not make λ_{db} and λ_{sb} be large simultaneously. Moreover, by the results in Sec. II, we know that Δm_{B_d} and $\sin 2\beta_{J/\Psi K_S}$ will push the allowed parameter space of λ_{db} to the region with small values. Without loss of generality, for simplicity we directly set the effects of λ_{db} be insignificant and ignorable. Hence, we will focus on the contributions of λ_{sb} in our numerical presentation, which relate to various $b \to s$ processes.

For numerical calculations and constraints, we list the useful values in Table I, where the relevant CKM matrix elements $V_{td} = |V_{td}| \exp(-i\beta_d)$ and $V_{ts} = -|V_{ts}| \exp(-i\beta_s)$ are obtained from the UTfit Collaboration [26], the decay constant of B_q is referred to the result given by the HPQCD Collaboration [27], the CDF and DØ average value of Δm_{B_s} is from Ref. [1] and the SM Wilson coefficients for $b \to q q' \bar{q}'$ and $b \to q\ell^+\ell^-$ are obtained from

Ref. [28]. The upper limit for $\mathcal{B}(B_s \to \mu^+\mu^-)$ with 95% confidence level (C. L.) is quoted by the latest DØ measurement [31]. Other inputs are from the particle data group (PDG) [25].

TABLE I. Experimental data and numerical inputs for the parameters in the SM.

$\lvert V_{td}\rvert$	β_d	$\lvert V_{ts}\rvert$	β_s	m_{B_d}
$8.51(22) \times 10^{-3}$	$(22 \pm 0.8)^\circ$	$4.07(22) \times 10^{-2}$	$-(1.03 \pm 0.06)^\circ$	5.28 GeV
m_{B_s}	$f_{B_d}\sqrt{\hat{B}_d}$	$f_{B_s}\sqrt{\hat{B}_s}$	f_{B_d}	f_{B_s}
5.37 GeV	(216 ± 15) MeV	(266 ± 18) MeV	190 ± 13 MeV	231 ± 15 MeV
$S^{\mathrm{Exp}}_{J/\Psi K_S}$	$S^{\mathrm{Exp}}_{\phi K_S}$	$(\Delta m_{B_d})^{\mathrm{Exp}}$	$(\Delta m_{B_s})^{\mathrm{Exp}}$	$\mathcal{B}^{\mathrm{Exp}}(b \to s\ell^+\ell^-)$
0.655 ± 0.024	$0.44^{+0.17}_{-0.18}$	0.507 ± 0.005 ps^{-1}	17.77 ± 0.12 ps^{-1}	$(4.5 \pm 1.0) \times 10^{-6}$
$\mathcal{B}^{\mathrm{Exp}}(B_s \to \mu^+\mu^-)$	C_3	C_4	C_5	C_6
$< 5.1 \times 10^{-8}$	0.013	-0.0335	0.0095	-0.0399
$C^{SM}_{7\gamma}$	C^{SM}_9	C^{SM}_{10}	$\sin^2\theta_W$	$\alpha(m_Z)$
-0.305	4.344	-4.430	0.231	1/129

Before we discuss the VQM predictions, it is necessary to know which processes involve less hadronic uncertainties and could give the strict constraints. We find that in addition to Δm_{B_s}, the observed inclusive $b \to s\ell^+\ell^-$ decays with $\ell = e, \mu$ are the good candidates. Although the possible constraint of $\sin 2\beta_s^{J/\Psi\phi}$ has been mentioned in Sec. II, as shown in Eq. (12), its current measurement cannot provide any significant bound. We present Δm_{B_s} (down-left hatched) and $\mathcal{B}(b \to s\ell^+\ell^+)$ (dotted) with 2σ errors of the data as functions of $|\lambda_{sb}|$ and θ^Z_s in Fig. 3, in which $|\lambda_{sb}|$ is in units of 10^{-3}. From the figure, we see that $\mathcal{B}(b \to s\ell^+\ell^-)$ further limits the upper value of $|\lambda_{sb}|$ to be around 10^{-3}. In general, the range of the CP violating phase θ^Z_s is $[-\pi, \pi]$. For simplicity, we just show the results within $[0, \pi]$. The pattern of the constraint in $[-\pi, 0]$ is similar to that in $[0, \pi]$.

Since we set the Z-mediated $b \to d$ transition be negligible, the wrong-sign charge asymmetry for B_d decays is ascribed to the SM contribution. We take $a^d_{s\ell}(SM) = -4.8 \times 10^{-4}$ for our numerical estimates. Using Eq. (26) for $a^s_{s\ell}$ and Eq. (4) for the like-sign charged asymmetry, the contours for $A^b_{s\ell}$ as a function of $|\lambda_{sb}|$ and θ^Z_s are shown in Fig. 4(a), where the numbers in the plot are units of 10^{-4}. We also plot $A^b_{s\ell}$ as a function of θ^Z_s with fixing $|\lambda_{sb}|$ in

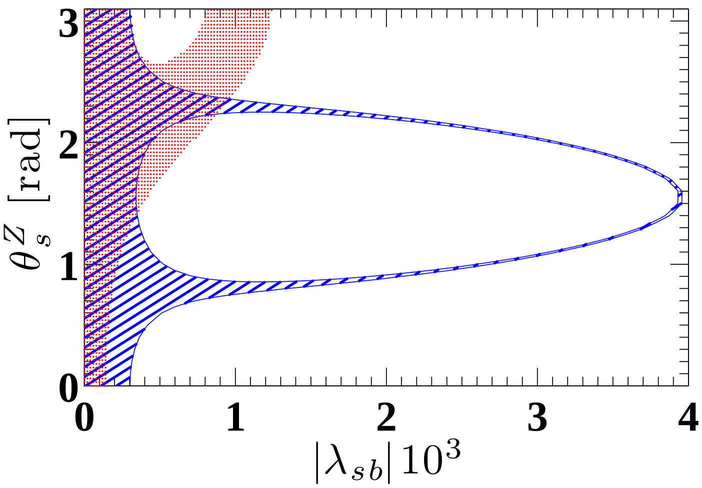

FIG. 3. Constraint of $|\lambda_{sb}|$ and θ_s^Z from Δm_{B_s} and $\mathcal{B}(b \to s\ell^+\ell^-)$.

Fig. 4(b), in which the solid, dashed and dash-dotted lines denote $|\lambda_{sb}| = (0.5, 0.7, 0.9)\times 10^{-3}$, respectively. From Fig. 4, we see that due to the constraint of $\mathcal{B}(b \to s\ell^+\ell^-)$, the absolute value of the like-sign charge asymmetry $A^b_{s\ell}$ can be as large as 5×10^{-4}. Although the result is not enhanced by order of magnitude, it could be still a factor of two larger than the SM prediction.

Next, we analyze the time-dependent CPA in the $B_s \to J/\Psi\phi$ decay. When Z-mediated $b \to d$ effects are neglected, it is easy to find that $A^b_{s\ell}$ and $S_{J/\Psi\phi}$ defined in Eq. (27) have a strong correlation. By using Eq. (34), the contours for the time-dependent CPA of $\sin 2\beta_s^{J\Psi\phi}$ as a function of $|\lambda_{sb}|$ and θ_s^Z are displayed in Fig. 5(a). Moreover, $\sin 2\beta_s^{J\Psi\phi}$ as a function of θ_s^Z with fixing $|\lambda_{sb}|$ is shown in Fig. 5(b) with the same legend as Fig. 4(b). According to the results, we find that when the constraints of Δm_{B_s} and $\mathcal{B}(b \to s\ell^+\ell^-)$ are taken into account at the same time, the sign of $\sin 2\beta_s^{J/\Psi\phi}$ favors negative, which is the same as that indicated by CDF and DØ measurements. Although the upper limit on the magnitude is smaller than the current data, it could still be 15%, whereas the SM prediction is only around 4%.

In terms of the early analysis, the penguin-like Z-mediated effect for $b \to sc\bar{c}$ in Eq. (30) could be estimated as

$$\left| \frac{\lambda_{sb} c_L^d}{V_{ts}^* V_{tb}} c_V^c \right| < 0.0048 \sim |a_5|. \tag{42}$$

It is clear that the new effect to the decay amplitude of $B \to J/\Psi K_S$ is insignificant as the case in the SM. Thus, we have $\sin 2\beta_{J/\Psi K_S} \approx \sin 2\beta_{J/\Psi K_S}(SM) \approx 0.695$ in the Z-mediated

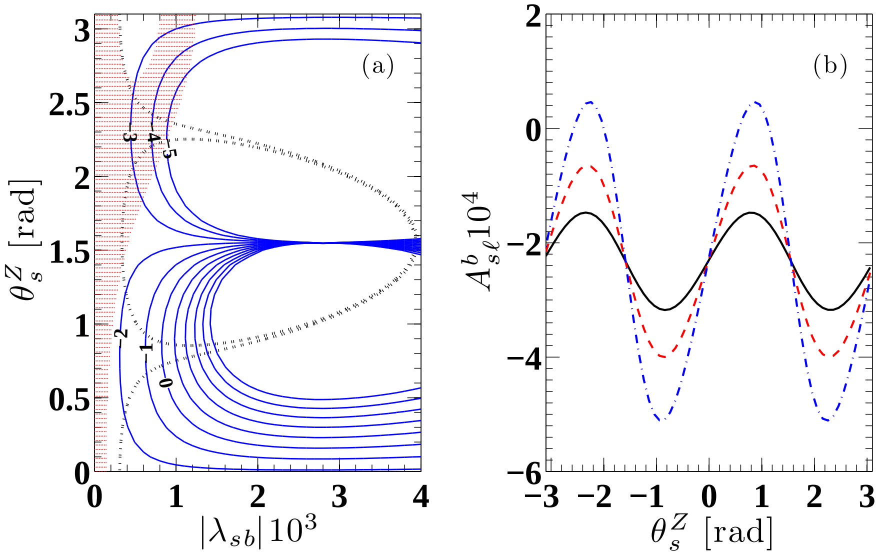

FIG. 4. (a) Contours for $A^b_{s\ell}$ (in units of 10^{-4}) as a function of $|\lambda_{sb}|$ and θ_s^Z and (b) $A^b_{s\ell}$ as a function of θ_s^Z, where the sold, dashed and dash-dotted lines represent $|\lambda_{sb}| = (0.5, 0.7, 0.9) \times 10^{-3}$, respectively.

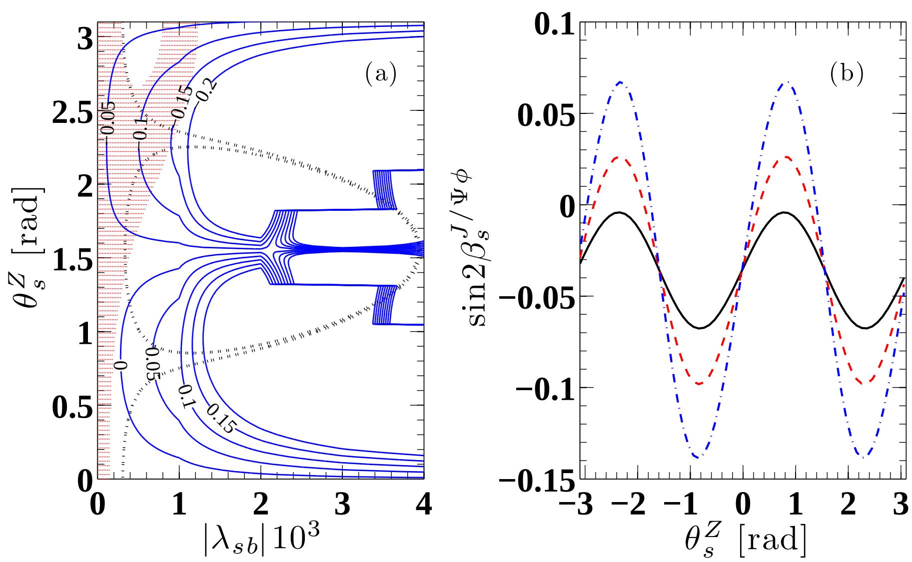

FIG. 5. (a) Contours for $\sin 2\beta_s^{J/\Psi\phi}$ as a function of $|\lambda_{sb}|$ and θ_s^Z and (b) $\sin 2\beta_s^{J/\Psi\phi}$ as a function of θ_s^Z with the same legend as Fig. 4(b).

VQM. In order to probe the new CP violating source arising from $SU(2)$ singlet exotic quarks, the best observable is the time-dependent CPA in the $B_d \to \phi K_S$ decay, where $\sin 2\beta_{\phi K_S}$ and $\sin 2\beta_{J/\Psi K_S}$, defined by Eqs. (33) and (34), have similar values in the SM, respectively. To understand the influence of Z-mediated effects on the CPA in $B_d \to \phi K_S$, we display the contours for $\sin 2\beta_{\phi K_S}$ as a function of $|\lambda_{sb}|$ and θ_s^Z in Fig. 6(a). From the result, we find that $\sin 2\beta_{\phi K_S}$ could approach 0.90 when $|A^b_{s\ell}|$ is a factor of two larger than the SM prediction. Furthermore, in Fig. 6(b) we present the contours for ΔS_{β_d}, the difference in the CPA between $J/\Psi K_S$ and ϕK_S modes defined by Eq. (35). Clearly, the difference of -20% could be achieved. It is interesting to mention that $\sin 2\beta_{\phi K_S}$ in the VQM is larger than $\sin 2\beta_{J/\Psi K_S}$ in $[0, \pi]$, whereas the situation is reversed in $[-\pi, 0]$. Although the current data in $B_d \to \phi K_S$ prefers the latter case, in this region $|A^b_{s\ell}|$ is even smaller than the SM result. Due to the current accuracy of the data, it is hard to tell which solution is more close to the reality. Hence, more precise measurements are necessary. For further comprehending the θ_s^Z dependence, we plot $\sin 2\beta_{\phi K_S}$ and ΔS_{β_d} as functions of θ_s^Z in Fig. 7(a) and (b), where the solid, dashed and dash-dotted lines denote $|\lambda_{sb}| = (0.5, 0.7, 0.9) \times 10^{-3}$, respectively.

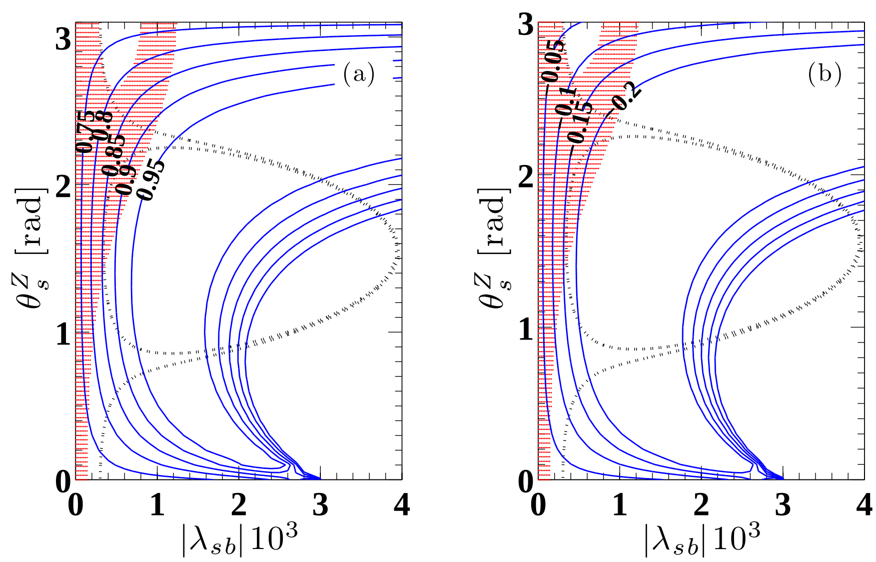

FIG. 6. (a) [(b)] Contours for $\sin 2\beta_{\phi K_S}$ [ΔS_{β_d}] as a function of $|\lambda_{sb}|$ and θ_s^Z.

Finally, we analyze the rare decays of $B_q \to \ell^+\ell^-$. As discussed earlier, the $b \to d$ transition in the Z-mediated VQM is suppressed and therefore, we will concentrate on

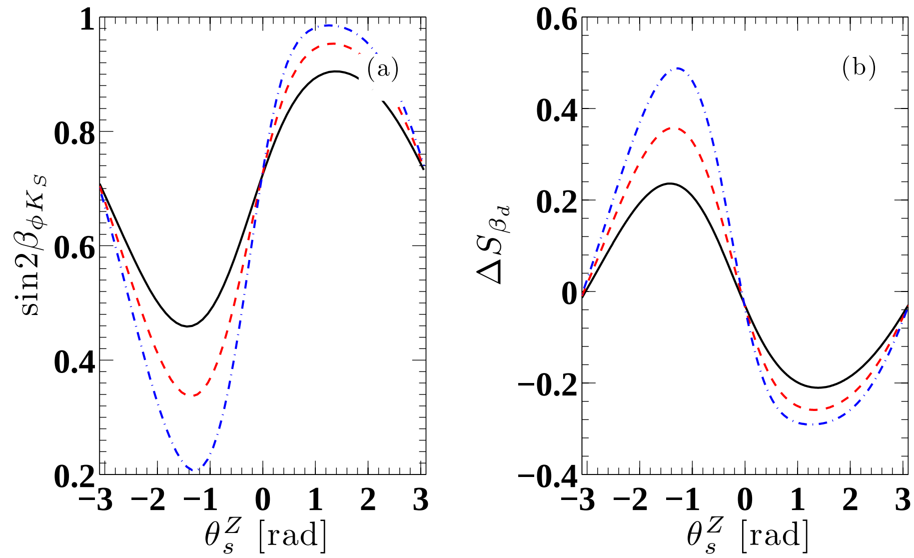

FIG. 7. (a) [(b)] $\sin 2\beta_{\phi K_S}$ [ΔS_{β_d}] as a function of θ_s^Z with the same legend as Fig. 4(b).

$B_s \to \ell^+\ell^-$. Since the leptonic process is helicity-suppressed, only the heavier charged leptonic modes are interesting. However, since the experiments only provide the limits on $B_s \to \mu^+\mu^-$, we study the influence of Z-mediated effects on the muon channel. Using Eq. (41) and the values in the Table I, the contours for $\mathcal{B}(B_s \to \mu^+\mu^-)$ as a function of $|\lambda_{sb}|$ and θ_s^Z are displayed in Fig. 8. We find that the upper value of $\mathcal{B}(B_s \to \mu^+\mu^-)$ is around 0.6×10^{-8} whereas the SM result of $\mathcal{B}^{SM}(B_s \to \mu^+\mu^-)$ is around 0.39×10^{-8}.

V. CONCLUSION

We have model-independently studied the charge and CP asymmetries as well as FCNCs in the various $B_{d,s}$ processes. In particular, we have found that $(-A^b_{s\ell}) < 3.16 \times 10^{-3}$ when the constraints from the $B_q - \bar{B}_q$ mixings and the time-dependent CP asymmetries (CPA) for $B_q \to J/\Psi M_q$ with $M_q = K, \phi$ and $q = d, s$ are taken into account. Although the upper value is smaller than the data of the new DØ measurement, it is still one order of magnitude larger than the standard model (SM) prediction and sensitive to new CP violating effects. We have also explored the VQM to illustrate the possible large effects on $|A^b_{s\ell}|$ and FCNCs in the $B_{d,s}$ processes. Explicitly, we have shown that (a) the like-sign charge asymmetry could be enhanced by a factor of two in magnitude; (b) the CPA of $\sin 2\beta_s^{J/\Psi\phi}$ could reach

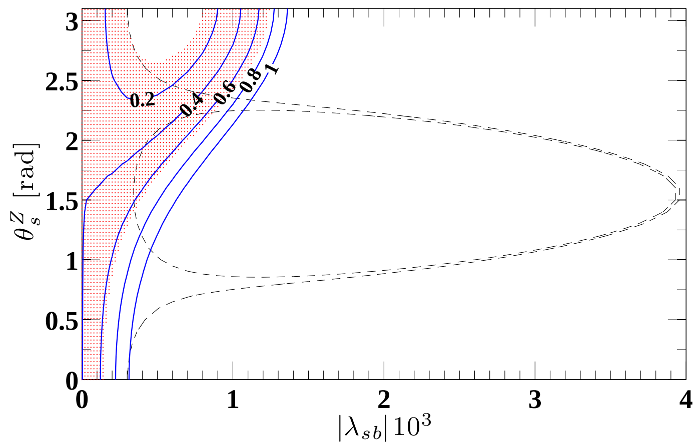

FIG. 8. Contours for $\mathcal{B}(B_s \to \mu^+\mu^-)$ as a function of $|\lambda_{sb}|$ and θ_s^Z, where the numbers in the plot are in units of 10^{-8}.

to -15%; (c) the CPA of $\sin 2\beta_{\phi K_S}$ could be higher than $\sin 2\beta_{J/\Psi K_S}$ when $|A^b_{s\ell}|$ is larger than the SM prediction; and (d) the BR for $B_s \to \mu^+\mu^-$ could be as large as 0.6×10^{-8}.

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